

# Dirac neutrinos and dark matter stability from lepton quarticity

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We propose to relate dark matter stability to the possible Dirac nature of neutrinos. The idea is illustrated in a simple scheme where small Dirac neutrino masses arise from a type-I seesaw mechanism as a result of a  $Z_4$  discrete lepton number symmetry. The latter implies the existence of a viable WIMP dark matter candidate, whose stability arises from the same symmetry which ensures the Diracness of neutrinos.

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## I. INTRODUCTION

Amongst the major shortcomings of the Standard Model are the neutrino mass and the dark matter problems. Underpinning the origin of neutrino mass and elucidating the nature of dark matter would constitute a gigantic step forward in particle physics. Here we focus on the possibility that the neutrino mass and dark matter problems may be closely interconnected [1, 2]. Concerning neutrino mass a major unknown is whether neutrinos are their own anti-particles, an issue which has remained an open challenge ever since Ettore Majorana's pioneer idea on the quantum mechanics of spin. On the other hand, since many years physicists have pondered about what is the dark matter made of, and what makes it stable, a property usually assumed in an *ad-hoc* fashion <sup>1</sup>. Indeed, although the existence of non-baryonic dark matter is well established by using cosmological and astrophysical probes, its nature has otherwise remained elusive [5].

The detection of neutrinoless double beta decay would be a major step in particle physics since, according to the black-box theorem [6, 7] it would not only demonstrate that lepton number is violated in nature, but also imply that neutrinos are of Majorana type. On the other hand, the fact that the weak interaction is V-A turns this quest into a major challenge [8, 9]. As of now the nature of neutrinos remains as mysterious as the mechanism responsible for generating their small masses. Little is known regarding the nature of its associated messenger particles, the underlying mass scale or its flavor structure [10], currently probed only in neutrino oscillation experiments [11]. Here we assume that neutrinos are Dirac particles. Realizing this possibility requires extra assumptions beyond the standard  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$  electroweak gauge invariance. One approach is to extend the gauge group itself, for example, by using the  $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$  gauge structure due to its special features [12]. In this framework it has recently been shown how to obtain a type-II seesaw mechanism for Dirac neutrinos [13, 14]. Using unconventional lepton charges for right handed neutrinos [15] and gauging  $B - L$  can also lead to Dirac neutrinos within the type-I seesaw mechanism [4, 16, 17]. Alternatively, one may stick to the  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$  gauge structure but use extra flavor symmetries implying a conserved lepton number, so as to obtain Dirac neutrinos, as suggested in [18].

In this letter we focus on having neutrinos as Dirac particles as a result of the cyclic symmetry  $Z_4$ , which plays the role of a discrete version of lepton number, we call *quarticity*. We show that a WIMP dark matter candidate can naturally emerge, stabilized by quarticity, the same symmetry associated to the Diracness of neutrinos. In Sect. II we sketch in some detail the required particle content of the model and demonstrate how the Dirac nature of

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<sup>1</sup> Attempts at stabilizing dark matter by using the  $Z_2$  and  $Z_3$  symmetries have already been considered in the literature [3, 4].

neutrinos and the stability of dark matter follow from the same principle. In Sect. III we discuss the main aspects concerning our WIMP dark matter candidate, including a brief discussion of the interactions relevant for determining its relic density. We also discuss its direct detection potential through the Higgs portal mechanism. Finally we summarize our results in Sect. IV.

## II. THE MODEL

Our model is based on the discrete symmetry  $Z_4 \otimes Z_2$  where  $Z_4$  is the cyclic group of order four and  $Z_2$  is the cyclic group of order two. The group  $Z_4$  can be viewed as a discrete version of lepton number. As we show here,  $Z_4$  symmetry not only forbids Majorana terms but also forbids couplings of potential dark matter candidate ensuring stability. Thus, the same symmetry which implies the Dirac nature of neutrinos also ensures the stability of dark matter. The  $Z_2$  symmetry is required here to ensure the seesaw origin of neutrino mass, by forbidding tree level coupling between the left and right handed neutrinos.

The particle content of our model along with the  $Z_4$  charge assignments of the particles are as shown in Table I.

Fields	$Z_4$	$Z_2$	Fields	$Z_4$	$Z_2$
$\bar{L}_{i,L}$	$\mathbf{z}^3$	$\mathbf{1}$	$\nu_{i,R}$	$\mathbf{z}$	$-\mathbf{1}$
$l_{i,R}$	$\mathbf{z}$	$\mathbf{1}$	$\bar{N}_{i,L}$	$\mathbf{z}^3$	$\mathbf{1}$
$N_{i,R}$	$\mathbf{z}$	$\mathbf{1}$			
$\Phi$	$\mathbf{1}$	$\mathbf{1}$	$\chi$	$\mathbf{1}$	$-\mathbf{1}$
$\zeta$	$\mathbf{z}$	$\mathbf{1}$	$\eta$	$\mathbf{z}^2$	$\mathbf{1}$

Table I. The  $Z_4$  and  $Z_2$  charge assignments for leptons and the scalars ( $\Phi$ ,  $\chi$ ,  $\zeta$  and  $\eta$ ). Here  $\mathbf{z}$  denotes the fourth root of unity, i.e.  $\mathbf{z}^4 = 1$ .

In Table I,  $L_{i,L} = (\nu_{i,L}, l_{i,L})^T$ ;  $i = e, \mu, \tau$  are the lepton doublets which have charge  $\mathbf{z}$  under  $Z_4$ . The  $l_{i,R}$ ;  $i = e, \mu, \tau$  are the charged lepton singlets which carry  $Z_4$ -charge  $\mathbf{z}$ . Apart from the Standard Model fermions, the model also includes three right-handed neutrinos  $\nu_{i,R}$  transforming as singlets under the  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$  gauge group and with charge  $\mathbf{z}$  under  $Z_4$ . We also add three gauge singlet Dirac fermions  $N_{i,L}, N_{i,R}$ ;  $i = 1, 2, 3$  with charge  $\mathbf{z}$  under  $Z_4$ , as shown in Table. I. All the fermions except  $\nu_{i,R}$  are even under the  $Z_2$  symmetry, while  $\nu_{i,R}$  are odd.

In the scalar sector the  $\Phi = (\phi^+, \phi^0)^T$  transforms as  $SU(2)_L$  doublet, does not carry any  $Z_4$  charge and is even under  $Z_2$  symmetry. On the other hand the real scalar  $\chi$  is a gauge singlet, carries no  $Z_4$  charge but is odd under  $Z_2$  symmetry. We also add two other gauge singlet scalars  $\zeta$  and  $\eta$  both of which carry  $Z_4$  charges  $\mathbf{z}$  and  $\mathbf{z}^2$  respectively and are even under  $Z_2$ . The fact that  $\mathbf{z}^2 = -1$  allows us to take the field  $\eta$  also to be real.

The  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes Z_4 \otimes Z_2$  invariant Yukawa Lagrangian for the leptons is given by

$$\begin{aligned}
-\mathcal{L}_{\text{Yuk,l}} = & y_{ij}^l \bar{L}_{i,L} \Phi l_{j,R} + f_{ij} \bar{L}_{i,L} \tilde{\Phi} N_{j,R} + g_{ij} \bar{N}_{i,L} \chi \nu_{j,R} + M_{ij} \bar{N}_{i,L} N_{j,R} \\
& + y_{ij}^\eta \bar{\nu}_{i,R}^c \nu_{j,R} \eta + \text{h.c.}
\end{aligned} \tag{1}$$

where  $y_{ij}, f_{ij}, g_{ij}, y_{ij}^\eta$ ,  $i, j = 1, 2, 3$ , are the Yukawa couplings. Also  $\tilde{\Phi}$  is the  $SU(2)_L$  conjugate field,  $\tilde{\Phi} = i\sigma_2 \Phi^*$  where  $\sigma_2$  is the second Pauli matrix. After symmetry breaking the  $Z_4$  neutral scalars  $\Phi, \chi$  acquire vacuum expectation values (vevs)

$$\langle \Phi \rangle = \frac{v}{\sqrt{2}} \quad \text{and} \quad \langle \chi \rangle = u.$$

The  $Z_4$  charged scalars  $\eta, \zeta$  acquire no vev thus ensuring that the  $Z_4$  symmetry remains unbroken even after electroweak symmetry breaking. Since the scalar  $\chi$  was odd under  $Z_2$  symmetry, its vev breaks the  $Z_2$  symmetry spontaneously. The neutrinos can then acquire a tiny mass through type-I Dirac seesaw mechanism as shown in Fig. 1 and briefly discussed below.

After symmetry breaking the charged leptons acquire mass through the vev of the  $SU(2)_L$  doublet scalar  $\Phi$ . The  $6 \times 6$  mass matrix for the neutrinos and heavy fermions  $N_{i,L}, N_{i,R}$  in the basis  $(\bar{\nu}_{i,L}, \bar{N}_{i,L})$  and  $(\nu_{i,R}, N_{i,R})^T$  is given by

$$M_{\nu,N} = \begin{pmatrix} 0 & f v / \sqrt{2} \\ g u & M \end{pmatrix}. \tag{2}$$

where  $f, g, M$  are each  $3 \times 3$  matrices with elements  $f_{ij}, g_{ij}, M_{ij}$  respectively. The elements of the invariant mass matrix  $M$  for the heavy leptons  $N_{i,L}, N_{i,R}$  can be naturally much larger than the symmetry breaking scales appearing in the off-diagonal blocks, i.e.  $M_{ij} \gg v, u$ . In this limit the mass matrix in (2) can be easily block diagonalized by the perturbative seesaw diagonalization method given in [19]. The resulting  $3 \times 3$  mass matrix for light neutrinos can be viewed as the Dirac version of the well known type-I seesaw mechanism and is given as follows

$$M_\nu = \frac{u v}{\sqrt{2}} f M^{-1} g \tag{3}$$

Note that  $\Phi$  and  $\chi$  are both singlets under  $Z_4$ . While  $\Phi$  is even under  $Z_2$ ,  $\chi$  is odd and its

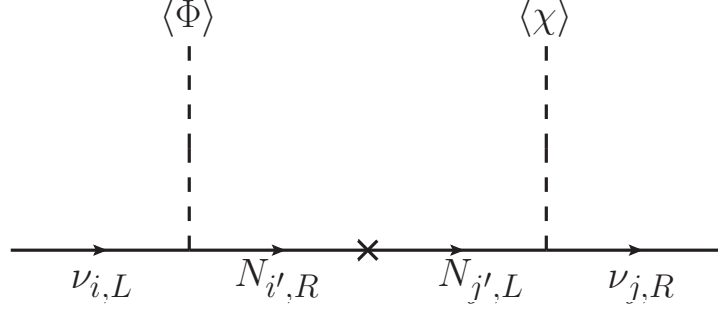


Figure 1. Dirac type-I seesaw mechanism.

vev breaks  $Z_2$  symmetry inducing neutrino mass.

Turning now to the  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes Z_4 \otimes Z_2$  invariant scalar potential, we can write it as

$$\begin{aligned}
 V = & -\mu_\Phi^2 \Phi^\dagger \Phi - \frac{\mu_\chi^2}{2} \chi^2 + \frac{m_\eta'^2}{2} \eta^2 + m_\zeta'^2 \zeta^* \zeta + \frac{\kappa}{2} \eta \zeta^2 + \frac{\lambda_\Phi}{4} (\Phi^\dagger \Phi)^2 + \frac{\lambda_\chi}{16} \chi^4 + \frac{\lambda_\eta}{16} \eta^4 \\
 & + \frac{\lambda_\zeta}{4} (\zeta^* \zeta)^2 + \frac{\lambda'_\zeta}{16} \zeta^4 + \frac{\lambda_{\Phi\chi}}{2} \Phi^\dagger \Phi \chi^2 + \frac{\lambda_{\Phi\eta}}{2} \Phi^\dagger \Phi \eta^2 + \lambda_{\Phi\zeta} \Phi^\dagger \Phi \zeta^* \zeta + \frac{\lambda_{\chi\eta}}{4} \chi^2 \eta^2 \\
 & + \frac{\lambda_{\chi\zeta}}{2} \chi^2 \zeta^* \zeta + \frac{\lambda_{\zeta\eta}}{2} \eta^2 \zeta^* \zeta + \text{h.c.}
 \end{aligned} \tag{4}$$

After symmetry breaking one has, in the unitary gauge

$$\Phi \rightarrow \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h' \end{pmatrix}, \quad \chi \rightarrow u + \chi' \tag{5}$$

The minimization conditions are given by

$$\begin{aligned}
 \mu_\Phi^2 &= \frac{\lambda_\Phi}{4} v^2 + \frac{\lambda_{\Phi\chi}}{2} u^2 \\
 \mu_\chi^2 &= \frac{\lambda_\chi}{4} u^2 + \frac{\lambda_{\Phi\chi}}{2} v^2
 \end{aligned} \tag{6}$$

The  $h'$  and  $\chi'$  fields mix with each other and their mass squared matrix in the  $(h', \chi')$  basis is given by

$$M_{h',\chi'}^2 = \frac{1}{4} \begin{pmatrix} \lambda_\Phi v^2 & 2\lambda_{\Phi\chi} u v \\ 2\lambda_{\Phi\chi} u v & \lambda_\chi u^2 \end{pmatrix}. \tag{7}$$

After diagonalization the CP-even neutral mass eigenstates  $h$  and  $\chi$  of 7 are given by

$$\begin{aligned}
 h &= \cos \theta h' + \sin \theta \chi' \\
 \chi &= -\sin \theta h' + \cos \theta \chi'
 \end{aligned} \tag{8}$$

where the doublet–singlet mixing angle  $\tan 2\theta = \frac{4\lambda_{\Phi\chi}uv}{\lambda_{\Phi}v^2 - \lambda_{\chi}u^2}$ . We identify  $h$  as the 125 GeV particle recently discovered at LHC.

Let us discuss briefly the role of the  $\zeta$  and  $\eta$  scalars which are singlets under the  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$  gauge group, are even under  $Z_2$  symmetry, but carry  $Z_4$  charges  $\mathbf{z}$  and  $\mathbf{z}^2$  respectively. Since  $\eta, \zeta$  fields do not acquire vevs, they do not mix with other fields. However, their masses do receive contribution from the vev of  $\Phi, \chi$  and are given by

$$\begin{aligned} m_{\eta}^2 &= m_{\eta}^{\prime 2} + \frac{\lambda_{\Phi\eta}}{2} v^2 + \frac{\lambda_{\chi\eta}}{2} u^2 \\ m_{\zeta}^2 &= m_{\zeta}^{\prime 2} + \frac{\lambda_{\Phi\zeta}}{2} v^2 + \frac{\lambda_{\chi\zeta}}{2} u^2 \end{aligned} \quad (9)$$

In the absence of  $\zeta$  and  $\eta$ , the Lagrangian of the model acquires an enhanced symmetry namely  $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y \otimes U(1) \otimes Z_2$  where  $U(1)$  is a continuous global symmetry which can be interpreted as a generalized global lepton number symmetry. In the presence of the scalars  $\zeta$  and  $\eta$  one can write following  $Z_4 \otimes Z_2$  invariant terms in the scalar potential

$$\eta^2, \eta\zeta^2, \eta^4, \zeta^4, \eta^2\zeta^*\zeta + \text{H.c.} \quad (10)$$

Notice that all these terms are  $Z_4 \otimes Z_2$  invariant but break the global  $U(1)$  invariance so the remaining symmetry group is just  $Z_4 \otimes Z_2$ . Note that apart from the above terms, the other terms in scalar potential 4 are invariant under the global  $U(1)$ . On the other hand the field  $\eta$  also couples to the right handed neutrinos through a  $Z_4 \otimes Z_2$  invariant term

$$\nu_{i,R}^T \sigma_2 \nu_{j,R} \eta + \text{h.c.} \quad (11)$$

Again this Yukawa coupling is only  $Z_4 \otimes Z_2$  invariant, breaking the continuous  $U(1)$  symmetry. Owing to the couplings of  $\eta$  to the scalar  $\zeta$  in 10 and to right handed neutrinos as in 11, the latter can also couple to  $\zeta$  as shown in Fig. 2.

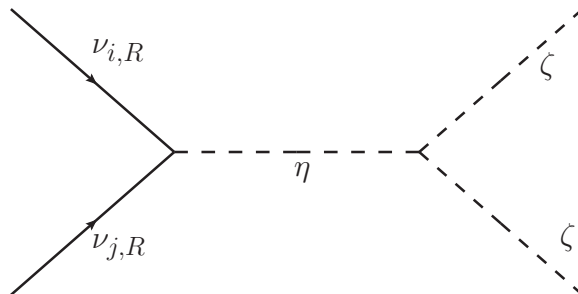


Figure 2. The coupling between right handed neutrinos and  $\zeta$  mediated by the scalar  $\eta$ .

After the scalars  $\Phi$  and  $\chi$  acquire vevs, the  $Z_2$  symmetry breaks spontaneously, implying a small mass for neutrinos. However, since neither  $\Phi$  nor  $\chi$  carries a  $Z_4$  charge, one has that

$Z_4$  remains exact even after spontaneous electroweak breaking. The unbroken  $Z_4$  symmetry ensures that the neutrinos do not acquire any Majorana mass term, retaining their Dirac nature even after spontaneous symmetry breaking takes place.

On the other hand, since the  $\zeta$  and  $\eta$  fields carry the  $Z_4$  charge, this implies that  $\zeta$  as well as  $\eta$  should not acquire any vev in order to prevent breaking  $Z_4$ . This ensures that neutrinos remain Dirac and also leads to the possibility that  $\zeta$  can be a stable particle and thus a potential candidate for dark matter. We discuss this possibility in the next section.

### III. WIMP SCALAR DARK MATTER CANDIDATE

We now turn briefly to the issue of stability of  $\zeta$  and its suitability as a dark matter candidate. As mentioned before, to ensure that neutrinos remain Dirac particles, the  $Z_4$  quarticity symmetry should remain exact, so that no scalar carrying a  $Z_4$  charge should acquire any vev. Thus both  $\eta$  and  $\zeta$  which carry  $Z_4$  charges can potentially be stable as a result of the unbroken  $Z_4$  symmetry. However, owing to the  $Z_4$  charge assignments, the  $\eta$  field has cubic couplings to both scalars and fermions as noted in Eqs. 10 and 11. These couplings lead to the decay of  $\eta$  and thus, despite carrying a  $Z_4$  charge,  $\eta$  is not stable.

On the other hand, due to its  $Z_4$  charge, all cubic couplings of  $\zeta$  are forbidden by the unbroken  $Z_4$  symmetry. This implies that there is no term of the form  $\zeta\rho_i\rho_j$  where  $\rho_i, \rho_j$  stands for other scalar species of the model which is allowed by  $Z_4$ . Likewise, in the Yukawa sector the  $Z_4$  symmetry also ensures that all terms of the form  $\zeta\psi_i\psi_j$ ,  $\psi_i$  denoting a generic fermion, are forbidden. Thus, the residual  $Z_4$  symmetry responsible for the Dirac nature of neutrinos also ensures the stability of the  $\zeta$  making it a potentially viable dark matter candidate.

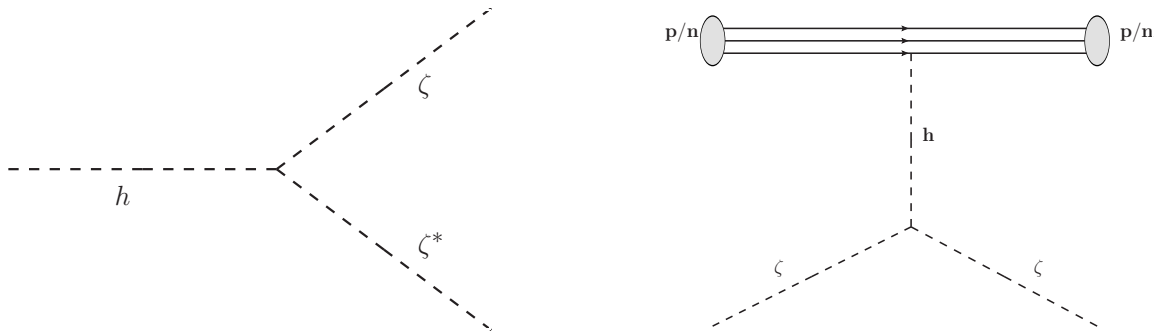


Figure 3. The diagrams for invisible Higgs decay to two dark matter particles and direct detection of dark matter through Higgs mediated nuclear recoil.

Although  $\zeta$  is stable, with no direct tree level coupling to fermions, owing to the symmetry of the model, it couples to right handed neutrinos through exchange of  $\eta$  as shown

in Fig. 2. Moreover, it interacts with other scalars through cubic and quartic terms of the type  $\zeta^*\zeta\rho$  and  $\zeta^*\zeta\rho_i^\dagger\rho_j$  respectively. In particular, the dark matter interaction with the Higgs  $h$  (we denote the recently discovered 125 GeV particle as “Higgs” which in our model will be an admixture of CP-even scalars given in 8) is of special phenomenological interest. If the dark matter mass  $m_\zeta < \frac{m_h}{2}$  then the decay  $h \rightarrow \zeta\zeta^*$ , shown diagrammatically in Fig. 3, is kinematically allowed. The decay width of  $h \rightarrow \zeta\zeta^*$  is given by

$$\Gamma(h \rightarrow \zeta\zeta^*) = \frac{\lambda_{h\zeta}^2 v^2}{16\pi m_h} \sqrt{1 - \frac{4m_\zeta^2}{m_h^2}} \quad (12)$$

where  $\lambda_{h\zeta}$  is the Higgs dark matter coupling constant. At LHC, this will be invisible decay of Higgs. Assuming no addition beyond SM decay apart from the invisible decay, the current LHC data puts a constraint on such invisible decay widths to be no more than 17% of the total decay width [20]. This puts a stringent constraint on the Higgs dark matter coupling  $\lambda_{h\zeta}$  as shown in Fig. 4. In our model, apart from the dark matter and other

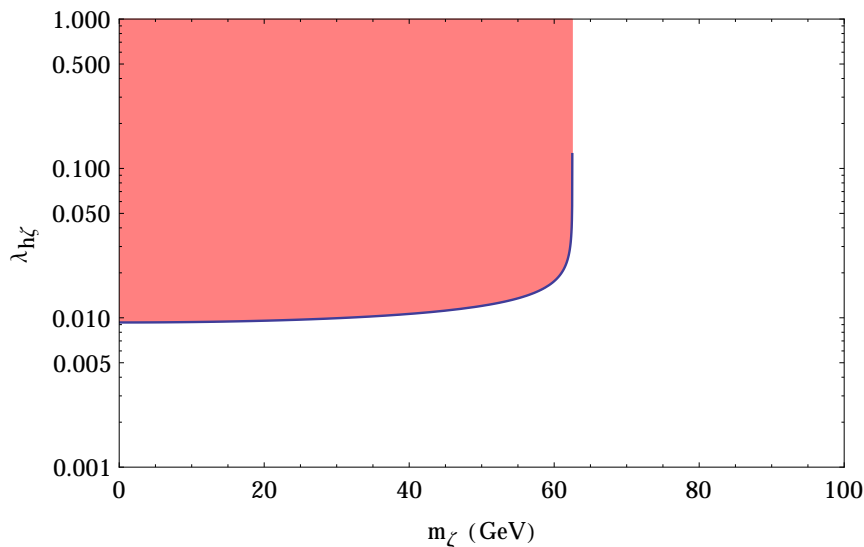


Figure 4. The constraint on Higgs dark matter coupling from invisible Higgs decay. The shaded region is ruled out from limits on Higgs invisible decay width [20].

Standard Model Higgs decay modes, one has decays to the new scalars  $\chi, \eta$ . In plotting Fig. 4, for simplicity, we have assumed that Higgs decay to  $\chi, \eta$  is kinematically forbidden i.e.  $m_\chi, m_\eta > \frac{m_h}{2}$ . It should be noted that if these decay modes are also allowed, they will also contribute to the invisible decay of Higgs as both  $\chi, \eta$  further decay only to neutrinos.

The Higgs dark matter coupling can also be used in order to detect dark matter directly through nuclear recoil in a variety of experiments [4, 21]. The process proceeds through tree-level Higgs exchange, as illustrated in Fig. 3. Thus our  $\zeta$  WIMP dark matter model



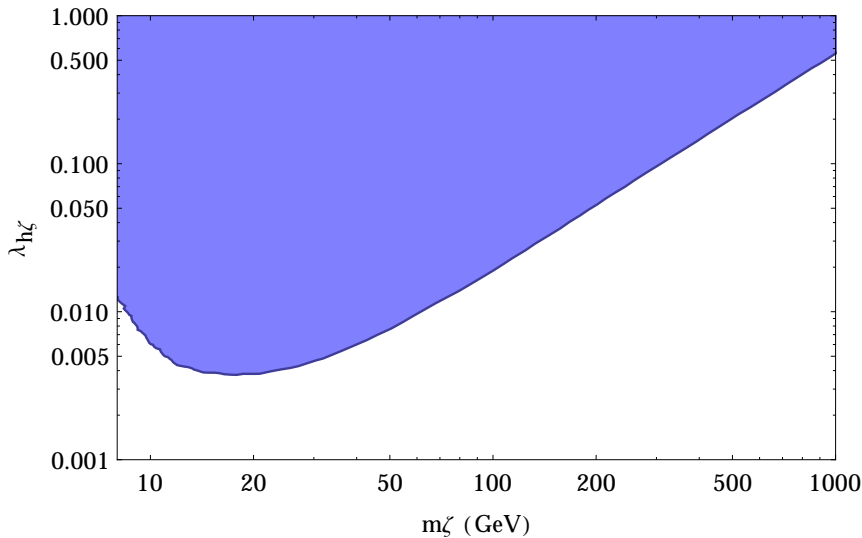


Figure 5. The experimental sensitivity to our WIMP scalar dark matter candidate. The shaded region is ruled out by LUX data [22].

realizes the so-called “Higgs portal” dark matter scenario. The Higgs-mediated nucleon–dark matter cross-section in our model is given by

$$\sigma = \frac{\lambda_{h\zeta}^2 f_n^2}{\pi m_h^4} \frac{m_n^4}{(m_n + m_\zeta)^2} \quad (13)$$

where  $m_n$  is the mass of the nucleon and  $f_n$  is the dimensionless Higgs–nucleon effective coupling constant. The null results from the direct detection experiments such as LUX [22] lead to rather stringent constraints on the coupling characterizing the interaction of our WIMP scalar dark matter with the Higgs boson. The experimental sensitivity to our WIMP scalar dark matter candidate is illustrated in Fig 5. In plotting Fig 5 we have taken the nucleon mass  $m_n = 0.93895$  GeV and the effective Higgs nucleon coupling constant  $f_n = 0.30$  [23]. One sees that, in view of the region excluded by the LUX experiment, constraints from the invisible Higgs decay  $h \rightarrow \zeta\zeta$  are relevant only in the low-mass region.

In our model we have also several terms that lead to dark matter annihilations into two other scalars and fermions, through various diagrams such as illustrated in Fig. 6. By solving numerically the relevant system of Boltzmann equations one can show that the model parameters can be chosen to lead to the correct relic density for dark matter [24] <sup>2</sup> In addition to the direct detection signal discussed here, these terms can also lead to indirect detection signatures associated, for example, to gamma-ray lines from the annihilation channels  $\zeta\zeta \rightarrow \gamma\gamma$  and  $\zeta\zeta \rightarrow Z\gamma$  channels [24]. A more detailed study of the discovery

<sup>2</sup> Dark matter transforming non-trivially under a given discrete symmetry has been previously studied in several works and shown indeed to provide a viable dark matter scenario [21, 23].

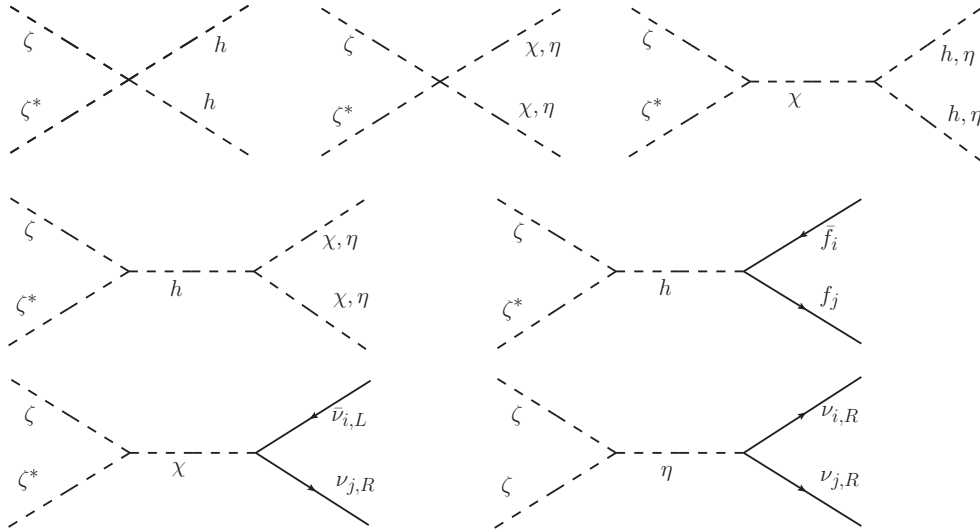


Figure 6. Typical diagrams contributing to the annihilation of WIMP dark matter

potential of our WIMP candidate in direct as well as indirect detection experiments will be presented elsewhere.

Before ending this section let us briefly mention other features of our model. First, in our model a conserved  $Z_4$  charge would also lead to the hypothetical lepton number violation quadruple beta decay process [25]. Moreover, since  $\eta$  is a real scalar field which couples to right handed neutrinos, its decay to two neutrinos or two antineutrinos may potentially generate a lepton asymmetry in the Universe, with potential application to leptogenesis with a conserved  $Z_4$  lepton number [26]. We intend to present details of these and other features and implications of our model in subsequent works. Before closing let us also mention that our model can easily be generalized by including vector-like quarks, so as to accommodate the recent diphoton hint seen by the ATLAS and CMS collaborations. It would be identified with the scalar  $\chi$ , very much along the lines of Refs. [27, 28].

#### IV. DISCUSSION AND SUMMARY

We have proposed that dark matter stability reflects the Dirac nature of neutrinos. We illustrated how to realise this proposal within the simplest type I seesaw mechanism for neutrino mass generation. The scheme naturally leads to a WIMP dark matter candidate which is made stable by the same discrete lepton number  $Z_4$  symmetry, or quarticity, which makes neutrinos to be Dirac particles. Dark matter can be probed by searching for nuclear recoil through the Higgs portal mechanism. The inclusion of the quark sector proceeds rather trivially. Notice that, in making our point that the Diracness of neutrinos may be responsible for dark matter stability, we have chosen a simple “flavor-blind” scheme based

on the use of the simplest  $Z_2$  symmetry to provide the seesaw mechanism. The later can be replaced by more complex symmetries, that could also tackle the problem of flavor and leptonic CP violation. We shall address this issue in a follow up work.

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*RS will like to dedicate this paper in memory of Dr. I. Sentitemsu Imsong.*

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